

# Flavor Alignment via Shining in RS

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## Abstract

We present a class of warped extra dimensional models whose flavor violating interactions are much suppressed compared to the usual anarchic case due to flavor alignment. Such suppression can be achieved in models where part of the global flavor symmetry is gauged in the bulk and broken in a controlled manner. We show that the bulk masses can be aligned with the down type Yukawa couplings by an appropriate choice of bulk flavon field representations and TeV brane dynamics. This alignment could reduce the flavor violating effects to levels which allow for a Kaluza-Klein scale as low as 2-3 TeV, making the model observable at the LHC. However, the up-type Yukawa couplings on the IR brane, which are bounded from below by recent bounds on CP violation in the  $D$  system, induce flavor misalignment radiatively. Off-diagonal down-type Yukawa couplings and kinetic mixings for the down quarks are both consequences of this effect. These radiative Yukawa corrections can be reduced by raising the flavon VEV on the IR brane (at the price of some moderate tuning), or by extending the Higgs sector. The flavor changing effects from the radiatively induced Yukawa mixing terms are at around the current upper experimental bounds. We also show the generic bounds on UV-brane induced flavor violating effects, and comment on possible additional flavor violations from bulk flavor gauge bosons and the bulk Yukawa scalars.

# 1 Introduction

Efforts to solve the hierarchy problem of the Standard Model (SM) usually end up introducing new physics at the TeV scale. However, most of these models would also allow new flavor violation effects which have not been observed. This usually leads one to require that the new physics is very nearly flavor universal, which would exclude the possibility of finding a testable solution to the fermion mass hierarchy problem within the same framework.

Warped extra dimensions [1] offer new approaches to flavor physics [2]. The puzzle of the hierarchical flavor structure can be solved via the split fermion mechanism [3], using flavor dependent wave functions for the Standard Model fermions [4]. This "anarchic" approach to flavor naturally yields the hierarchies of masses and mixing angles. Within this framework, the same setup also induces extra protection against excess of flavor changing neutral current (FCNC) processes [5,6] (see also [7]), as there is a built-in RS-GIM mechanism [6] where flavor violation is dominantly due to non-universality of fermion profiles near the IR brane. Thus, flavor and CP violation (CPV) related to the first two generations is highly suppressed [6,8] as in the SM case. A residual little CP problem is, however, still found in the form of too large contributions to the neutron electric dipole moment (EDM) [6], and sizable chirally enhanced contributions to  $\epsilon_K$  [9–13]. FCNCs are dominated by the exchange of Kaluza-Klein (KK) gluons, which induce the most dangerous (LR-type)  $\Delta F = 2$  four-Fermi operators contributing to  $K - \bar{K}$  mixing. The lower bound on the KK mass scale (taking into account the effect of the RS-GIM mechanism) is, at leading order, around 20 TeV [9] (compared with  $\sim 10^4 - 10^5$  TeV [12] for flavor-structureless models), which would render the model unobservable at the LHC. Recently in [14] it was argued, based on matching the full RS set-up onto a two site model, that if the Higgs is in the bulk and one loop matching of the gauge coupling is included, the KK scale can be lowered down to around 5-6 TeV. The weaker bound is due to a combination of several effects: (i) one loop matching of the gauge coupling which lowers the leading order bound by a factor of two, down to roughly 10 TeV, above the LHC reach [15]; (ii) exploiting effects related to the presence of a bulk Higgs which amounts to effectively increasing the overall scale of the 5D Yukawa coupling, allowing for more elementary light fermions. In the analysis below, we shall apply the one loop matching result, item (i). Increasing the size of the 5D Yukawa however, item (ii), implies that RS chirality violating processes such as the ones which contribute to EDMs,  $b \rightarrow s\gamma$  [6] are enhanced. In particular, it was recently shown that the bound from  $\epsilon'/\epsilon_K$  [16] constrains the KK scale to be above the  $\mathcal{O}(6)$  TeV scale, still above the LHC reach (see also [17] for Higgs mediated flavor violation processes which are also enhanced in this limit). Thus, in our study we assume an IR brane localized Higgs, and our results can be viewed as conservative ones in the sense that various constraints might be further ameliorated with a bulk Higgs.

The little CP problem might be accidentally solved if a combination of various *unrelated* CP violating parameters just happen to be small, of  $\mathcal{O}(0.1)$ . However, in view of the fact that the CKM phase is of  $\mathcal{O}(1)$  we are motivated to look for a parametric solution for the problem. One may seek guidance by using the AdS/CFT correspondence [18], where the RS setup can be understood from a 4D point of view [19], identifying 4D global currents with

bulk gauge fields. In fact, in the context of electroweak precision tests, progress was made [20] by gauging the approximate custodial symmetry of the SM in the bulk, as motivated by the duality. Thus, we expect that a similar progress can be made by gauging (part or all of) the SM approximate flavor symmetries [21–29].

There are several directions of how to realize models with gauged 5D flavor symmetries. One possible distinction is whether flavor issues (the flavor puzzle and the flavor problem) are addressed solely by Planck physics on the UV brane or whether bulk physics participates in the flavor dynamics as well. Obviously, our ability to directly and indirectly probe flavor dynamics in the near future depends on which of these cases is realized in nature. In [21,22,29] it was proposed to break the flavor symmetries only on the UV brane. Such a setup leads to a class of models where the flavor puzzle is solved by Planck scale physics and described by 4D general minimal flavor violation (GMFV) [30,31] (where the 4D Yukawa matrices control the flavor violation but generically one needs to resum the contributions related to the third generation fermions [31]). This class of models is in general similar in structure to the SM, hence hard to probe via flavor precision tests.

Another class of models which was recently proposed [23,24,28] (see also [25–27] for the lepton sector) try to solve the little CP problem but without giving up on addressing the flavor puzzle (and taking advantage on the built-in RS-GIM). These models may allow us to directly probe new flavor dynamics at the the LHC, which is very exciting, as well as leading to deviate from the SM predictions in indirect flavor precision tests. The basic idea is to align the down type quark sector which includes the bulk masses and the 5D down Yukawas (and possibly also the brane kinetic terms) such that the constraint from  $\epsilon_K$  is satisfied. The model proposed in [24] is based on combining two U(1) horizontal symmetries broken on the UV and IR branes respectively and an additional flavor diagonal U(1) symmetry (motivated by the custodial symmetry to  $Z \rightarrow b\bar{b}$  [32]) which allows one to split-up the SM weak doublet partners into two separate 5D multiplets. In [28] a global U(3) symmetry for the down quark weak singlet fields was proposed in order to provide the required alignment.<sup>1</sup> In this model one has to assume a mild hierarchy in the 5D down Yukawa parameters, and there is also a tension related to mixing with KK states due to the absence of RS-GIM protection for the RH down type quarks (see [25] for a related discussion). In [23], 5D MFV (5DMFV) was proposed as a framework which may allow to control the form of flavor violation in warped models. Generic 5DMFV models at low energies yield a flavor structure which is similar to generic RS models, and hence do not solve the little CP problem [23]. However, 5DMFV allows to ameliorate the problem via accidental approximate alignment: a consequence of a single flavor diagonal parameter,  $r_u$ , which controls the tree level down quark flavor violation, being accidentally small. As discussed below, however, the RS  $\epsilon_K$  problem requires a rather high degree of alignment,  $r_u = \mathcal{O}(10^{-2})$ , which calls for a non-accidental solution.

In this work we examine how to combine aspects of 5DMFV and alignment in order to construct a viable class of models in which the assumptions on the UV (Planckian) breaking of flavor are minimal. An important tool which we use here is the mechanism

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<sup>1</sup>Since it is expected that the bulk physics, which include quantum gravity effects, would generically violate any global symmetries, we imagine that this symmetry is actually gauged.

of “shining” [33], *i.e.* transmitting a symmetry breaking effect from the UV brane through the bulk by scalar fields. This idea was first applied to warped space flavor models in [21] to explain how natural flavor conservation can appear in an RS type model. In [21], the fermions were assumed to be confined to the IR brane, so in that case one could protect the fermions from flavor violation but there was no explanation for the flavor hierarchy. The main new ingredient here is that we are also allowing the field which transmits the flavor violation to couple to the bulk fermions, resulting in small flavor violating effects in the bulk via higher dimensional operators. These effects are then manifested as non universal bulk masses for the quarks, giving rise to flavor dependent wave functions for the quark zero modes, like in [23], thus explaining the flavor hierarchy.

We will present two different symmetry breaking patterns: In the first model, we impose a  $SU(3)_Q \times SU(3)_d$  flavor symmetry, while in the second one we impose only the diagonal subgroup  $SU(3)_{Q+d}$ . In both cases we assign the bulk scalars with the quantum numbers of the down-type Yukawa coupling  $Y_d$ , namely  $(\mathbf{3}, \bar{\mathbf{3}})$  in the first case, and  $\mathbf{8}$  in the second. We do not impose any symmetry in the right-handed up sector, thus allowing for anarchic bulk masses, but we require that the up-type Yukawa coupling is an IR brane field, so that it does not feed into VEVs of other bulk fields. This approach actually leads to the fact that the up type sector tends to be anarchic, hence (as in other alignment models - see *e.g.* [23, 24, 34] and references therein), the presence of up type flavor and CPV is expected in our scenario. Recently, however, it was pointed out [35] that a strong bound from CPV in  $D^0 - \bar{D}^0$  mixing already exists and is translated to a strict lower bound on the scale of the 5D up type Yukawa matrix in anarchic RS models. This implies that the tree-level alignment is going to be spoiled due to higher order effects induced by the sizable up type Yukawa matrix. We will examine these misaligning effects in great detail, and show possible ways of reducing their effects below the current experimental bounds. We also comment on possible flavor effects of the bulk scalars and flavor gauge bosons.

The paper is organized as follows: in section 2 we give a basic realization of 5DMFV, and show that in order to have alignment, one needs a more restricted symmetry breaking pattern than the generic 5DMFV of [23]. In section 3, we present the two improved models in which the down type Yukawa is aligned with the bulk quark masses. In section 4 we discuss misalignment effects beyond the leading order. These include effects at the IR brane due to the presence of the up-type Yukawa background field, effects of both IR and UV brane kinetic terms. We comment on possible FCNCs from flavons and flavor gauge fields. We conclude in section 5.

## 2 The Setup

First we construct a model where bulk scalars transmit (shine) flavor violation through the bulk, giving a concrete realization to 5DMFV. We then show that the suppression of flavor violating effects in generic 5DMFV models is quite a bit less than originally expected, motivating us to look for models with complete alignment.

## 2.1 Generic 5DMFV from Shining without Alignment

We start with some definitions and notations, which will be used throughout this paper. Recall that using conformally flat coordinates, the RS metric is given by a slice of  $\text{AdS}_5$

$$ds^2 = \left(\frac{R}{z}\right)^2 (dx_\mu dx_\nu \eta^{\mu\nu} - dz^2), \quad (2.1)$$

bounded by a UV brane at  $z = R$  ( $R$  is also the AdS curvature scale) and an IR brane at  $z = R'$ . The magnitudes of the scales are given by  $R^{-1} \sim M_{Pl}$  and  $R'^{-1} = 1 - 2 \text{ TeV}$ .

The electroweak gauge group is extended to  $\text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)_{B-L}$  gauge symmetry in the bulk to incorporate custodial symmetry [20]. This symmetry is reduced on the UV brane to the SM group  $\text{SU}(2)_L \times \text{U}(1)_Y$ . We assign the following representations for the quarks:

$$Q_i(2, 1, 1/3), \quad u_i(1, 2, 1/3), \quad d_i(1, 2, 1/3) \quad (i = 1, 2, 3), \quad (2.2)$$

so that there are two separate  $\text{SU}(2)_R$  doublets per generation.

We impose the flavor symmetry  $\text{SU}(3)_Q \times \text{SU}(3)_u \times \text{SU}(3)_d$  in the bulk. The improved models discussed in section 3 have only a subgroup of this as their bulk global symmetry. Gauging this group will give rise to flavor gauge bosons which mediate FCNCs once the symmetry is broken. We will also assume that the only sources for flavor violation are on the branes. In the general case, these will break the flavor symmetry in the UV completely (we will later show the bounds on UV flavor violating effects). The UV flavor violation is then transmitted (shined) through the bulk via bulk scalars which couple to the UV flavor sources and acquire VEVs.

In order to transmit the flavor violation, we introduce bulk scalar fields  $y_{u,d}$  which transform under the flavor group (the precise representation of the bulk fields is model dependent and is discussed below). These scalars will be the origin of the brane localized Yukawa couplings, so that the 5D Yukawa terms then become

$$\lambda_u y_{u,ij} \bar{Q}_i \tilde{H} u_j + \lambda_d y_{d,ij} \bar{Q}_i H d_j, \quad (2.3)$$

These scalars  $y_{u,d}$  are assumed to have small bulk masses compared to the AdS curvature, and to first approximation their expectation values are uniform (it is easy to extend the discussion to non-constant VEVs, see the appendix). For example, consider a simple case, where the bulk fields,  $y_{u,d}$ , are bi-fundamental of the 5D flavor group, so that in our notation the usual dimensionless Yukawas are  $Y_{u,d} = \lambda_{u,d} \langle y_{u,d} \rangle R^{3/2}$ . The main difference between this approach and that of [21] is that we are also considering the effects of small higher dimension operators coupling the bulk fermions to the flavor symmetry breaking bulk scalars. This coupling will be the source of the flavor hierarchy. Although the coefficients of these operators are suppressed by the cutoff scale, their effect can still be very significant, due to the exponential sensitivity of the effective 4D masses on the bulk mass parameters.

## Effective Bulk Fermion Mass

To order  $Y^3$ , the effective bulk fermion masses (in units of the AdS curvature) are of the form

$$c_Q = \alpha_Q \cdot 1 + \beta_Q Y_d Y_d^\dagger + \gamma_Q Y_u Y_u^\dagger \quad (2.4)$$

$$c_u = \alpha_u \cdot 1 + \gamma_u Y_u^\dagger Y_u \quad (2.5)$$

$$c_d = \alpha_d \cdot 1 + \beta_d Y_d^\dagger Y_d. \quad (2.6)$$

Here the  $\alpha_i$  are flavor-invariant bulk mass parameters which are  $\mathcal{O}(1)$ . These can not be forbidden by any flavor symmetry. Although this term was omitted in [23], it does play an important role both for producing the leading order operator (compared to which the insertions of the bulk scalars are considered as perturbations), and for getting the desired alignment in flavor space.

The parameters  $\beta_i, \gamma_i$  are coefficients of higher dimension operators suppressed by the cutoff scale  $\Lambda$ . Let us estimate how large  $\beta, \gamma$  could be. The corresponding effective operators are of the form *e.g.*  $B y_u^\dagger y_u \bar{Q} Q$ , where  $B$  is a coefficient of dimension  $-2$ . Naive Dimensional Analysis (NDA) then suggests that  $B \lesssim (4\pi/\Lambda)^2$ , leading to  $\beta \equiv B/R^2 \lesssim (4\pi/\Lambda R)^2$ , which may be a significant correction to the leading term. In order to get a reliable prediction for the bulk masses  $c_i$ , one would also expect that the correction to the universal term after the insertion of the VEVs is small, so that the next correction is truly negligible numerically. However, the improved models we will present in the next section exhibit complete alignment in the bulk, so that while a higher order term would change the numerical values of the expansion coefficients, it would not spoil the alignment itself.

## Flavor Changing Neutral Currents

Given Eq.(2.4), the 4D mass matrices are given by

$$\begin{aligned} m_{ij}^{(u)} &= \frac{v}{\sqrt{2}} F_Q Y_u F_u, \\ m_{ij}^{(d)} &= \frac{v}{\sqrt{2}} F_Q Y_d F_d, \end{aligned} \quad (2.7)$$

where we have used the normalized IR values of the quark zero-mode wave functions,  $f_{Q,u,d}$  [6], as being the eigenvalues of the spurions  $F_{Q,u,d}$  (in the basis where the bulk masses,  $c_{Q,u,d}$ , are diagonal)

$$f_x = \sqrt{\frac{1 - 2c_x}{1 - (\frac{R}{R'})^{1-2c_x}}}, \quad (2.8)$$

Let us show explicitly the result of [23] that in generic GMFV models tree-level FCNC's are similar in rate to the anarchic case. For concreteness we focus on the down-type quark sector (since these processes are the ones most severely constrained, though the analysis is completely analogous for the FCNC's involving up-type quarks). We use the bulk  $U(3)^3$  flavor

symmetry to diagonalize the matrix  $Y_d$ , which is achieved by a bulk basis transformation of the  $Q_L$  and the  $d_R$ . We can perform a further bulk unitary transformation on  $u_R$ , after which the form of the bulk Yukawas is

$$Y_u = V_5 U^{(diag)}, \quad Y_d = D^{(diag)}. \quad (2.9)$$

Note that since  $Y_u$  feeds into the bulk mass parameter of  $Q_L$ , diagonalizing  $Y_d$  is *not* sufficient to diagonalize the down-type mass matrix. In fact, at this point,  $c_Q$  is not diagonal (while  $c_{u,d}$  are). The  $c_Q$  can be diagonalized by  $Q_L \rightarrow U_Q Q_L$ . Now one can calculate the zero mode wave functions for these bulk fermions, and the effective down-type brane Yukawa coupling takes the form

$$\bar{Q}_L i f_{Q_i} U_{Q_{ij}}^\dagger D_{jk}^{(diag)} f_{dk} d_{Rk} \quad (2.10)$$

To diagonalize this matrix we need an additional rotation on the 4D zero mode fields  $Q_L \rightarrow V_Q Q_L$ ,  $d_R \rightarrow V_d d_R$ , such that  $v \times V_Q^\dagger F_Q U_Q^\dagger D^{(diag)} F_d V_d = \text{diag}(m_d, m_s, m_b)$ . The origin of FCNC's is then the following: with diagonal bulk mass terms the fermion couplings to neutral gauge bosons is diagonal, but not universal, due to the different wave functions of the fermions. Then rotating the non-universal diagonal coupling matrices to the actual 4D mass basis will induce flavor changing couplings. The leading effect comes from the first KK mode of the gluon,

$$g_{s*} \left[ \bar{Q}_L V_Q^\dagger F_Q \gamma_\mu \gamma(c_Q) F_Q V_Q Q_L + \bar{d}_R V_d^\dagger F_d \gamma_\mu \gamma(c_d) F_d V_d d_R \right] G^{(1)\mu}, \quad (2.11)$$

where  $\gamma(c) \approx \frac{\sqrt{2}x_1}{J_1(x_1)} \frac{0.7}{6-4c}$ , and  $g_{s*}$  is the bulk gauge coupling in units of the AdS curvature. Therefore, indeed the generic 5DMFV framework contains four-Fermi operators of the form  $\bar{d}_{L,R} \gamma_\mu d_{L,R} \bar{d}_{L,R} \gamma^\mu d_{L,R}$ . The most dangerous of these is the LLRR, or  $(V-A)(V+A)$  operator, which is absent from the SM and has to be suppressed which requires some form of alignment to make the model viable [23].

## 2.2 Approximate Alignment

Let us now consider the limiting cases when the left handed bulk mass only couples to either up or down type flavor:

$$c_Q = \alpha_Q \cdot 1 + r_u \beta_Q Y_u Y_u^\dagger + r_d \gamma_Q Y_d Y_d^\dagger \quad (2.12)$$

in the limit of  $r_u \rightarrow 0$  or  $r_d \rightarrow 0$ . Consider the case  $r_u = 0$ . In this case we choose a bulk basis where  $Y_d$  is diagonal, so that the bulk masses for the  $Q_L$  fields and for  $d_R$  are automatically diagonal. Then the 4D mass matrix for the down type quarks is automatically diagonal in the same basis where the couplings to the bulk gauge fields are diagonal, so that there are no tree-level FCNCs involving down-type quarks. The mass matrix for the up-type quarks is not diagonal, and the misalignment between the two directions in flavor space implies the presence of up-type quark FCNCs. Since in this case the bulk mass matrices are aligned with the down-type Yukawa coupling, flavor violations occur only in the up

sector. The case analyzed in [23] corresponds to  $r_u \ll r_d$ , where down-type quark FCNCs are naively suppressed by  $r_u^2$ . However, this is not always the case: the suppression also depends on the values of the 5D CKM mixing angles, and usually the actual suppression is quite a bit less than the naive  $r_u^2$ . The origin of this effect can be understood as follows: due to the 5D GIM mechanism the universal part in  $c_Q$  has no flavor content and it cannot induce large CP breaking required to produce the CKM order one phase (this is just a manifestation of the 5D Kobayashi-Maskawa mechanism). On the same lines the suppression in FCNCs is proportional to the alignment of the *traceless* parts of  $c_Q$  and  $Y_d^\dagger Y_d$ . The actual suppression therefore depends on the amount of deviation from degeneracy between the various eigenvalues of the  $r_u$  and  $r_d$  terms in Eq. (2.12). An explicit example for this which can be treated analytically is the two flavor case, which we present below. Then we proceed to a numerical analysis of the suppression in the three flavor case.

## FCNC Suppression in Two Generations

In the two flavor case, the 5D Yukawa couplings can be parameterized as

$$Y_u = V_5 Y_u^{\text{diag}} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} y_{u,1} & 0 \\ 0 & y_{u,2} \end{pmatrix}, \quad (2.13)$$

$$Y_d = \begin{pmatrix} y_{d,1} & 0 \\ 0 & y_{d,2} \end{pmatrix}, \quad (2.14)$$

where  $V_5 = V_5(\theta)$  is the two flavor 5D CKM matrix. Just like in the SM with two flavors, CP is conserved and the Yukawas are real. The bulk mass  $c_Q$  is given by

$$c_Q = \alpha_Q \mathbf{1} + \beta_d^L Y_d Y_d^\dagger + \beta_u^L Y_u Y_u^\dagger, \quad (2.15)$$

$$= \alpha_Q \mathbf{1} + \beta_d^L \text{diag}(y_{d,1}^2, y_{d,2}^2) + \beta_u^L V_5 \text{diag}(y_{u,1}^2, y_{u,2}^2) V_5^T. \quad (2.16)$$

In this basis,  $c_d$  and  $c_u$  are already diagonal. Off-diagonal elements of  $c_Q$  signal tree-level FCNCs in the down-type quark sector. Let us introduce the matrix  $U_Q = U_Q(\phi)$  which diagonalizes  $c_Q$ . The angle  $\phi$  parametrizes the misalignment between  $c_Q$  and the  $Y_d$ , and measures the size of FCNCs in the down sector. We can find  $U_Q$  by solving

$$\begin{aligned} U_Q^T \{ \text{diag}(c_Q^1, c_Q^2) \} U_Q &\stackrel{!}{=} c_Q - \frac{1}{2} \text{tr}[c_Q] \\ &= \frac{1}{2} \beta_d^L (y_{d,1}^2 - y_{d,2}^2) \sigma_3 + \frac{1}{2} \beta_u^L (y_{u,1}^2 - y_{u,2}^2) V_5 \sigma_3 V_5^T, \end{aligned} \quad (2.17)$$

where we have included only the traceless contributions to  $c_Q$  since terms proportional to the identity do not affect  $U_Q(\phi)$ . In the limit of small misalignment ( $\sin \phi \ll 1$ ) we find that

$$\sin \phi = \kappa \cos \theta \sin \theta + \mathcal{O}(\kappa^2), \quad (2.18)$$

where

$$\kappa \equiv \frac{\beta_u^L (y_u^1)^2 - (y_u^2)^2}{\beta_d^L (y_d^1)^2 - (y_d^2)^2} \quad (2.19)$$



is the parameter relevant for the suppression of tree-level FCNCs, assuming no fine-tuning of the 5D CKM angle  $\theta$ .

As expected we find that tree-level FCNCs are suppressed if the contribution of the traceless part of  $Y_u Y_u^\dagger$  to  $c_Q$  is small compared with that of the traceless part of  $Y_d Y_d^\dagger$ .

### FCNC Suppression in Three Generations

We can use the analytical insights of the previous section to understand the more complicated three-flavor case. Unfortunately, no simple formula is available which relates the misalignment to the coefficients and Yukawa eigenvalues. However, we have learned that  $r = \beta_u^L/\beta_d^L$  alone is not a good predictor for small FCNCs. Certainly in the limit  $r \rightarrow 0$ , FCNCs vanish. Due to the other factor involving the ratio of traceless contributions, the suppression can be weaker than  $\sim r^2$ . We will now argue that in fact we generically expect less suppression because the ratio of traceless contributions to  $Y_u Y_u^\dagger$  and  $Y_d Y_d^\dagger$  is usually larger than one. The size of the respective traceless contributions is determined by the size of the non-degeneracy of  $c_u$  and  $c_d$ . Due to the hierarchies in the CKM and the mass spectrum [36], where

$$\frac{m_d}{m_s \lambda_C} \sim 0.2, \quad \frac{m_d}{m_b \lambda_C^3} \sim 0.09, \quad \frac{m_u}{m_c \lambda_C} \sim 0.009, \quad \frac{m_u}{m_t \lambda_C^3} \sim 0.0006, \quad (2.20)$$

and also since the down quark singlet masses are in the (exponentially suppressed) region where  $c_d > 0.5$  (see *e.g.* [6] for a detailed discussion) the degeneracy in  $c_d$  is much stronger than in  $c_u$  (see section 3). We therefore generically expect the traceless contribution to  $Y_u Y_u^\dagger$  to be larger than the one to  $Y_d Y_d^\dagger$  which in consequence enhances the misalignment between  $c_Q$  and  $Y_d$ . We have confirmed this argument numerically. In Fig.1, we show the dependence of  $|C_4^{5D-MFV}/C_4^{RS}|$  on the 5D CKM angles for a reference model with  $c_u = (0.68, 0.53, -0.06)$ ,  $c_d = (0.65, 0.6, 0.58)$ ,  $c_Q = (0.64, 0.59, 0.46)$ . We have fixed the ratio  $r = \beta_u^L/\beta_d^L = 1/4$  and scanned over the 5D CKM angles. We only plot points which result in the correct 4D CKM matrix and which do not have singular 5D CKM matrices where 5D Kobayashi-Maskawa mechanism would render a CP conserving model (this is the reason for the lack of points around  $\pi/2, \pi, \dots$ ). We clearly see that most of the points give much less suppression than the naive expectation  $r^2 \sim 1/16$ . Assuming no tuning in the 5D CKM matrix we conclude that only the limit  $r \rightarrow 0$  seems to sufficiently suppress dangerous FCNCs. This points towards a symmetry which we explore in the next section.

## 3 Complete Alignment via Shining at Leading Order

We have seen above that a suppression of the  $Y_u Y_u^\dagger$  term in the expression for  $c_Q$  is generically not sufficient to suppress down-type FCNCs. It seems that the only simple solution to this problem is to identify possible symmetry breaking patterns which automatically give  $r_u = 0$ . In this case the alignment is complete, irrespective of the 5D mixing angles.

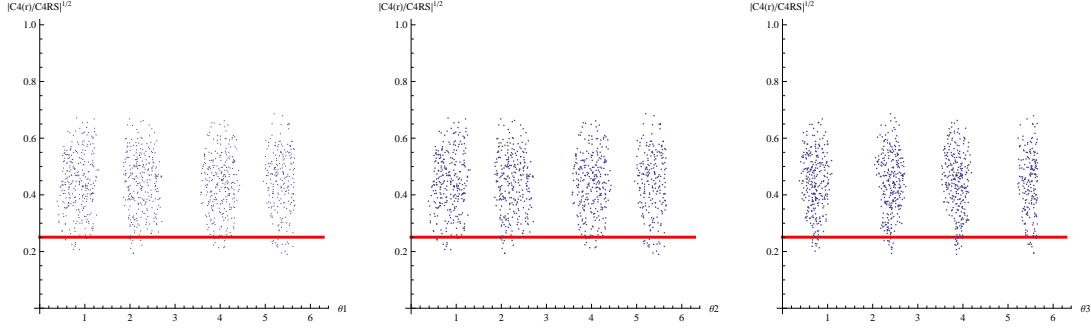


Figure 1: We show the ratio of Wilson coefficients  $|C_4^{5D-MFV}/C_4^{RS}|^{\frac{1}{2}}$  vs. the naive expectation  $r = 0.25$  (red line) as a function of the 5D CKM angles  $\theta_1, \theta_2$ , and  $\theta_3$ . We included only points that result in approximately the correct CKM matrix and Jarlskog determinant.

Here we consider two different models. In the first model, we impose  $SU(3)_Q \times SU(3)_d$  bulk flavor symmetry, while in the second one we impose only the diagonal subgroup  $SU(3)_{Q+d}$ . In both cases, the bulk scalar transforms as the down-type Yukawa  $Y_d$  itself, namely  $\mathbf{3}_Q \times \bar{\mathbf{3}}_d$  or  $\mathbf{8}_{Q+d}$  (assuming an adjoint field), respectively. We do not impose any flavor symmetries for the up sector, so that it is anarchic. Such scenario can be realized by having a flavor symmetry which is broken in the bulk by multiple sources but not via flavon fields which transform under the  $\mathbf{3}_Q \times \bar{\mathbf{3}}_d$  group such as a bulk  $Y_u$  field. The up-type Yukawa coupling is generated only on the IR brane, so that it cannot feed into the bulk fermion masses  $c_{Q,d}$ . In these models the alignment in the bulk between  $c_{Q,d}$  and  $Y_d$  is complete at leading order. Higher order effects are discussed in the next section.

### 3.1 $SU(3)_Q \times SU(3)_d$ Bulk Symmetry

In the first realistic example we assume that the down sector transforms under a  $SU(3)_Q \times SU(3)_d$  bulk flavor symmetry and that the only breaking is given by  $Y_d(\mathbf{3}_Q, \bar{\mathbf{3}}_d)$ . Then the expression for the bulk mass parameters is expected to be of the form

$$c_Q = \alpha_Q \cdot \mathbb{1} + \beta_Q Y_d Y_d^\dagger \quad (3.1)$$

$$c_d = \alpha_d \cdot \mathbb{1} + \beta_d Y_d^\dagger Y_d \quad (3.2)$$

$$m_{ij}^{(d)} = \frac{v}{\sqrt{2}} f_{Q_i} Y_{dij} f_{d_j}. \quad (3.3)$$

Note that in the above we have omitted, for simplicity, higher order terms. These terms would not spoil the alignment, but once resummed, they would only lead to a shift in the value of the eigenvalues [31]. In the RH up-sector we assume a generic breaking of the flavor symmetries by an anarchic Hermitian matrix.

We can find a numerically acceptable solution in the following way:

1. In the basis where  $Y_d$  is diagonal, we fix the input values of the  $c_Q$  matrix by requiring that it reproduces the physical CKM matrix. In particular, we used three sets of  $c_Q$  eigenvalues: the standard choice

$$c_Q = (0.632, 0.585, 0.430), \quad (3.4)$$

another one with a more composite third generation

$$c_Q = (0.610, 0.561, 0.200), \quad (3.5)$$

and finally close to fully composite third generation

$$c_Q = (0.602, 0.552, 0.000). \quad (3.6)$$

2. Eq. (3.1) then fixes  $Y_d$  as a function of  $\alpha_Q$  and  $\beta_Q$ .
3. Using (3.2), we obtain an expression for  $c_d$  as a function of the four input parameters  $\alpha_{Q,d}, \beta_{Q,d}$ .
4. We then scan over these four parameters, requiring that Eq.(3.3) reproduces the observed SM down-type quark masses.

Numerical examples resulting in the correct CKM and masses for the down quarks are given in Table 1.

$c_{Q_3}$	$\alpha_Q$	$\beta_Q \cdot 10$	$\alpha_d$	$\beta_d \cdot 10^2$	$Y_d$
0.000	1.258	-0.239	0.774	-0.228	(5.237, 5.433, -7.252)
0.000	-1.186	0.318	0.528	0.324	(7.498, 7.392, -6.107)
0.200	0.635	-0.733	0.629	-0.263	(0.591, 1.009, 2.437)
0.430	0.697	-0.944	0.635	-1.985	(0.833, 1.092, 1.683)
0.430	0.754	-0.636	0.655	-1.316	(1.388, 1.633, 2.258)

Table 1: Numerical examples for the model with  $SU(3)_Q \times SU(3)_d$  bulk symmetry resulting in a better than 95% C.L. fit of the mixing angles and masses.

We can see that in these solutions the Yukawa couplings are not hierarchical, and they are really implementations of the anarchic approach with additional bulk alignment. One might be concerned about the fact that in these numerical solutions, the correction to the flavor universal term is not always small (while perturbativity is always maintained). This would imply that the next correction might significantly modify the actual numerical values of the coefficients. However, as already mentioned, these higher order terms would not spoil the alignment itself, and therefore the absence of down-type tree-level FCNC is maintained.

### 3.2 $SU(3)_{Q+d}$ Bulk Symmetry

The second model that we are considering employs a diagonal flavor symmetry  $SU(3)_{Q+d}$  in the bulk, which is broken only by an adjoint spurion  $A_d(\mathbf{8}_{Q+d})$ . The main difference compared with the  $SU(3)_Q \times SU(3)_d$  case is that the alignment between  $c_{Q,d}$  and the Yukawa  $Y_d$  is not quadratic but linear, resulting in a very different numerical solution. The expressions for the bulk masses consistent with these symmetries are given by

$$c_Q = \alpha_Q \cdot \mathbb{1} + \beta_Q A_d \quad (3.7)$$

$$c_d = \alpha_d \cdot \mathbb{1} + \beta_d A_d \quad (3.8)$$

$$m_{ij}^{(d)} = \frac{v}{\sqrt{2}} f_{Q_i} (\alpha_Y \cdot \mathbb{1} + \beta_Y A_d)_{ij} f_{d_j} \quad (3.9)$$

where  $c_u$  is again assumed to be unconstrained. The procedure of finding numerical solutions for this system is slightly easier:

1. Again, we go to a basis where  $c_{Q,d}$  and  $A_d$  are diagonal, and fix the same input  $c_Q$  as before.
2. This fully fixes the direction of the adjoint matrix  $A_d$  in the space of  $\mathbb{1}, \lambda_3, \lambda_8$ , where  $\lambda_{3,8}$  are the diagonal Gell-Mann matrices. Writing it as  $c_Q = a_0 \cdot \mathbb{1} + a_3 \cdot \lambda_3 + a_8 \cdot \lambda_8$ , we find for example that  $a_0 = 0.55, a_3 = 0.033, a_8 = 0.14$  for the standard choice of  $c_Q$ .
3. The expression for  $c_d$  can then be written as  $c_d = \alpha_d \cdot \mathbb{1} + \beta_d \cdot (\lambda_3 + \frac{a_8}{a_3} \lambda_8)$ .
4. Now we can obtain  $f_{c_d}$  as function of  $\alpha_d, \beta_d$ , from which one can calculate the alignment of the  $m^d$  mass matrix in the  $\lambda_3, \lambda_8$  space.
5. Requiring that  $m^d$  points in the same direction in this space fixes the value of  $\alpha_d$  for any input value of  $\beta_d$ . This is performed numerically.

A few simple numerical examples are given in Table 2. We again find that the perturbative treatment of Yukawa insertions makes sense and that we do not need hierarchies in the Yukawa matrices to generate the SM CKM and mass spectrum.

## 4 Misalignment from loop effects

In the above, we have shown that, at leading order, we can completely eliminate the down-type FCNCs from bulk effects by aligning  $c_Q$  and  $c_d$  with  $Y_d$ . Since we assume that the up-type Yukawa (which breaks the flavor symmetries needed for the alignment) is an IR brane field, the alignment in the bulk could be complete. However, on the IR brane there will be higher dimension operators and loop effects which feed  $Y_u$  into  $Y_d$ , so that the brane induced  $Y_d$  is misaligned with the bulk  $c_{Q,d}$ . These effects are still placing very significant bounds on the model, some of which could actually be comparable to the tree-level bounds in the generic anarchic model found in [9]. The fact that some NLO effects could be very

$c_{Q_3}$	$\alpha_Q$	$\beta_Q$	$\alpha_d$	$\beta_d \times 10$	$\alpha_Y$	$\beta_Y$	$A_d$	$Y_d$
0.000	0.384	1.000	0.700	1.13	7.926	0.204	(0.217, 0.167, -0.384)	(7.984, 7.970, 7.825)
0.000	0.384	1.000	0.685	0.998	5.930	0.957	(0.217, 0.167, -0.384)	(6.137, 6.090, 5.562)
0.200	0.457	0.100	0.631	0.189	1.428	0.204	(1.529, 1.039, -2.568)	(1.740, 1.640, 0.904)
0.200	0.457	0.100	0.601	0.0276	0.679	-0.233	(1.529, 1.039, -2.568)	(0.323, 0.437, 1.278)
0.430	0.549	0.100	0.580	0.135	0.722	-0.299	(0.830, 0.359, -1.188)	(0.474, 0.615, 1.077)
0.430	0.549	0.100	0.590	0.195	0.823	-0.281	(0.830, 0.359, -1.188)	(0.589, 0.722, 1.158)

Table 2: Numerical examples for the model with  $SU(3)_{Q+d}$  bulk symmetry resulting in a better than 95% C.L. fit of the mixing angles and masses.

strong is not completely unexpected, since in order to suppress the tree-level CP violation in the  $D$  system [35] and to obtain a sufficiently large top mass, one usually needs to push the Yukawas close to their perturbative bounds, which in turn enhances the ratio of loop effects and also increases the importance of higher dimensional operators.

First we present the NDA estimates of all the couplings relevant for generating misalignments via loops. Then we present the bounds on the parameters from loop- (and higher dimensional operators-) induced effects: the first is contributing to the down Yukawa scalars on the IR brane; the second is kinetic mixing for the down-type quarks on the IR brane; the third is kinetic mixing on the UV brane; finally we comment on all other possible sources of additional flavor violation.

## 4.1 NDA for everything

The theory we are considering has an intrinsic cutoff scale  $\Lambda$ , which implies that the non-renormalizable theory becomes strongly coupled at that scale. In order to have a regime where the theory behaves as a weakly coupled 5D theory, we require that at least the first few KK modes are weakly coupled. We denote the number of weakly coupled KK modes by  $N_{\text{KK}}$  and usually require  $N_{\text{KK}} \geq 3$ , which implies  $\Lambda \sim N_{\text{KK}}/R$  for the unwarped cutoff scale.

### Brane localized Yukawa coupling

As a warm-up exercise, let us first consider the NDA bound for a brane Yukawa coupling of the form

$$\int d^4x \, y_d \bar{Q} H d \Big|_{z=R'}. \quad (4.1)$$

Here  $y_d$  is a dimension  $-1$  coupling. The NDA bound is obtained by requiring that the one-loop contribution in Fig. 2 to this coupling is not larger than the tree-level term itself.

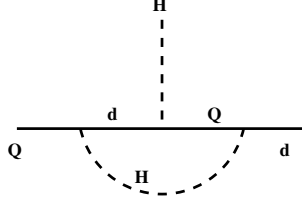


Figure 2: A one-loop diagram for corrections to a brane Yukawa coupling.

The estimate for the one-loop correction is given by

$$\Delta y_d^{1-loop} \sim \frac{4y_d^3}{16\pi^2} \Lambda^2 \quad (4.2)$$

where we have taken into account that the KK modes in the loop are coupled by a factor of  $\sqrt{2}$  more strongly than zero modes. The requirement that this does not exceed the tree-level bound will give the upper bound on the dimensionless Yukawa coupling  $Y_d = y_d R^{-1}$

$$Y_d < \frac{2\pi}{N_{\text{KK}}} \quad (4.3)$$

which is the usual bound on the dimensionless Yukawa couplings usually imposed on generic RS flavor models.

### Brane Yukawas from flavor scalars

What we are interested in here is the bound on the coefficients of brane localized Yukawa-type operators which originate from insertions of bulk or brane scalar fields breaking the flavor symmetries. As suggested in the previous section, we will assume that the down-type flavor symmetry is broken by a bulk scalar  $y_d$ , while the up-type by a brane scalar  $y_u$ . We are then interested in the NDA estimate of the size of the operator involving the  $y_u$  field in a down-type Yukawa coupling, which will give rise to misalignment and FCNC's. So we want NDA bounds on the coefficients  $\lambda_d$ ,  $\lambda_u$  and  $\lambda_d^u$  for the operators

$$S_{IR} \ni \int d^4x \left[ \lambda_d y_d \bar{Q} d H + \lambda_u y_u \bar{Q} u H + \lambda_d^u y_u y_u^\dagger y_d \bar{Q} d H \right] \Big|_{z=R'} \quad (4.4)$$

Since we assume that  $y_d$  is a bulk field and  $y_u$  is a brane field the coefficients  $\lambda_d$ ,  $\lambda_u$ ,  $\lambda_d^u$  are of mass dimension  $E^{-5/2, -2, -9/2}$  respectively. The leading corrections (see Fig. 3) arise at two loops for  $\lambda_d$  and  $\lambda_u$ , while for  $\lambda_d^u$  at four loops. Therefore (using the same rules as before) we find the leading corrections to be:

$$\Delta \lambda_d = \frac{8\lambda_d^3}{(16\pi^2)^2} \Lambda^5, \quad \Delta \lambda_u = \frac{4\lambda_u^3}{(16\pi^2)^2} \Lambda^4, \quad \Delta \lambda_d^u = \frac{8(\lambda_d^u)^3}{(16\pi^2)^4} \Lambda^9. \quad (4.5)$$

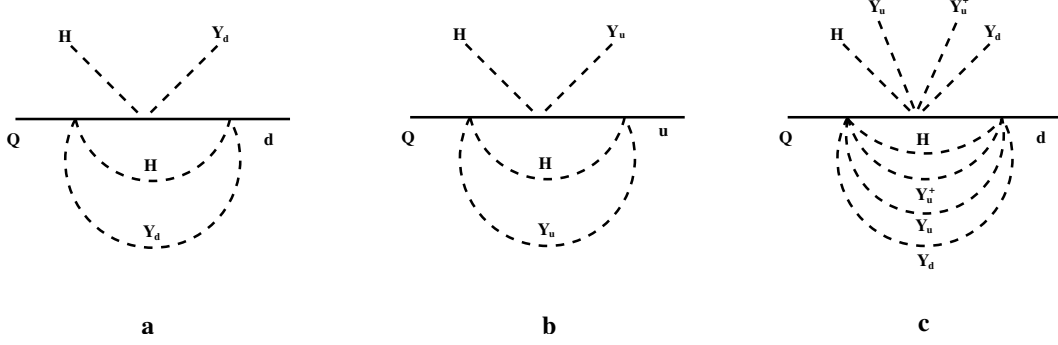


Figure 3: The leading corrections to the coefficients  $\lambda_d$  in **a**,  $\lambda_u$  in **b** and  $\lambda_d^u$  in **c** used estimating the NDA values of these operators.

The resulting NDA bounds will thus be

$$\lambda_d < \frac{16\pi^2 R^{\frac{5}{2}}}{\sqrt{8}N_{\text{KK}}^{\frac{5}{2}}}, \quad \lambda_u < \frac{16\pi^2 R^2}{2N_{\text{KK}}^2}, \quad \lambda_d^u < \frac{(16\pi^2)^2 R^{\frac{9}{2}}}{\sqrt{8}N_{\text{KK}}^{\frac{9}{2}}}. \quad (4.6)$$

As a consistency check, one can switch off the value of  $\lambda_d^u$  and regenerate it at one loop via the  $\lambda_{d,u}$  couplings. The contribution is generated at one-loop and is proportional to  $4\lambda_u^2 \lambda_d N_{\text{KK}}^2 / (R^2 16\pi^2)$ . Substituting the NDA values of  $\lambda_{u,d}$  in this expression, we indeed find the NDA value of  $\lambda_d^u$ .

### Scalar VEVs

In order to find the NDA value of the effective dimensionless Yukawa couplings  $Y_{u,d}$  in our model, we have to find (besides the coefficients  $\lambda_{u,d}$ ) also the NDA values of the bulk and brane scalar VEV's  $\langle y_{u,d} \rangle$ . For this we assume that there is a brane mass of order the cutoff scale generated, and in addition a brane localized quartic self-interaction. For the down-type bulk scalar  $y_d$

$$V(y_d)_{IR} = \Lambda y_d^\dagger y_d + \lambda_4 |y_d|^4, \quad (4.7)$$

where  $\lambda_4$  is of mass dimension  $-2$ . The NDA value of  $\lambda_4$  is obtained in a similar way to the above derivations: there is a one loop quadratically divergent correction

$$\Delta\lambda_4 = 4\lambda_4^2 \frac{\Lambda^2}{16\pi^2} \quad (4.8)$$

Hence the NDA value of  $\lambda_4$  is

$$\lambda_4 \sim \frac{4\pi^2 R^2}{N_{\text{KK}}^2}. \quad (4.9)$$

Thus the NDA value of the scalar VEV  $\langle y_d \rangle$  is

$$\langle y_d \rangle \sim \frac{\Lambda^{\frac{1}{2}}}{\lambda_4^{\frac{1}{2}}} = \frac{N_{\text{KK}}^{3/2}}{2\pi R^{3/2}}. \quad (4.10)$$

We also note that since  $y_u$  is an IR field, it follows the ordinary 4D power counting, and hence

$$\langle y_u \rangle \sim \frac{N_{\text{KK}}}{4\pi R}. \quad (4.11)$$

We have two important remarks regarding the NDA values of the VEVs:

- While all previous NDA estimates were *upper bounds* requiring that couplings don't blow up before the cutoff scale, finding the NDA value of the VEV  $\langle y_{d,u} \rangle$  involves minimizing a potential assuming the maximal NDA size quartic coupling. The VEV would actually *increase* with a decreasing quartic self-coupling. Thus for VEVs the NDA size is *not* an upper bound, merely the natural value. If one allows some cancellation between the tree-level and loop induced quartic, then the VEV can be *increased* without leaving the region of validity of the effective theory, but at the price of a tuning. For example, if there is a 50% cancelation in  $\lambda_4$  the VEV can be increased by a factor of  $\sqrt{2}$ .
- One does expect the size of the VEV to be set by the local infrared potential, even if the VEV started out with a different size on the UV brane. The reason is that one usually finds that by solving the bulk equations for a scalar VEV one usually finds that the VEV will "twist" between the UV and the IR brane values, ie. interpolate between the two natural NDA values set on the branes.

### Effective Yukawas from scalar VEVs

The NDA value effective Yukawa coupling for the bulk scalar is

$$Y_d = \lambda_d \langle y_d \rangle R^{-1} = \frac{2\sqrt{2}\pi}{N_{\text{KK}}}. \quad (4.12)$$

This is a factor of  $\sqrt{2}$  larger than the NDA bound for a brane localized Yukawa, Eq. (4.3), whose origin can be traced back to the fact that in the NDA bound for  $\lambda_d$  there is one brane localized field running in the loop which does not pick up an enhanced coupling. In addition, we have explained above that raising the VEV  $\langle y_d \rangle$  does not imply leaving the domain of validity of the effective theory, but rather a certain amount of tuning between the tree-level and loop induced quartics for  $y_d$  on the IR brane. For example, to reach  $Y_d = 10$  with  $N_{\text{KK}} = 3$  one needs a tuning of about 10%.

Similarly, if  $Y_u$  is generated from a brane scalar as proposed before, its NDA value is

$$Y_u = \lambda_u \langle y_u \rangle R^{-1} = \frac{2\pi}{N_{\text{KK}}}. \quad (4.13)$$



## 4.2 Loop induced corrections to the down Yukawa coupling

We are now able to estimate the misaligning effects of the IR brane loop corrections (or similarly the ones from higher dimensional operators). We have already estimated the NDA size of the  $\lambda_d^u$  operators in Eq.(4.6). This teaches us that the effective down-type Yukawa  $Y_d$  is shifted to:

$$Y_d \rightarrow Y_d^{(tot)} = Y_d + \lambda_d^u \langle y_u \rangle \langle y_u^\dagger \rangle \langle y_d \rangle R^{-1} = Y_d + \left( \frac{\lambda_d^u}{\lambda_d \lambda_u^2} \right) Y_u Y_u^\dagger Y_d R^2. \quad (4.14)$$

Using the NDA values of the  $\lambda$ 's, we find

$$Y_d^{(tot)} \sim Y_d \left( 1 + \frac{N_{KK}^2}{4\pi^2} Y_u Y_u^\dagger \right), \quad (4.15)$$

which is generated by the 1-loop diagram in Fig. 4. So the suppression parameter for the misaligning effects is given by

$$\epsilon_u = \frac{Y_u^2 N_{KK}^2}{4\pi^2} \quad (4.16)$$

where  $Y_u$  is an average value of the elements of the  $Y_u$  matrix. Note that this diagram is generically present in any alignment model where the Yukawas are dynamical fields, and operators of this sort cannot be forbidden by any symmetry. Since our theory is just on the verge of being perturbative, such corrections can be very dangerous. In fact, a recent analysis [35] shows that bounds on CPV in  $D - \bar{D}$  mixing requires  $Y_u > 1.6$ , and obtaining a sufficiently heavy top mass would also suggest a large value of  $Y_u$  ( $Y_u = 0.5$  is the lowest possible value when both the LH and RH top fields are fully composite). This implies that the suppression parameter  $\epsilon_u$  is numerically not that small. For  $Y_u = 1.6$  and  $N_{KK} = 3$  we find  $\epsilon_u = 0.58$  (while for  $N_{KK} = 2$  we would get  $\epsilon_u = 0.26$ ).

Due to the non-trivial flavor structure of (4.15) one needs to analyze carefully the down-type mass matrix in order to actually read off what the suppression of the flavor violating KK gluon couplings are. An easy way to do this is to evaluate the misalignment (in the basis where both the wave functions  $f_{Q,u,d}$  and  $Y_d$  are diagonal) from the left:

$$\begin{aligned} A_Q &\equiv f_Q Y_d^{(tot)} f_d^2 \left( Y_d^{(tot)} \right)^\dagger f_Q \\ &\sim f_Q Y_d f_d^2 Y_d f_Q + \frac{4N_{KK}^2}{16\pi^2} f_Q [Y_d f_d^2 Y_d Y_u Y_u^\dagger + Y_u Y_u^\dagger Y_d f_d^2 Y_d] f_Q \\ &\quad + \left( \frac{4N_{KK}^2}{16\pi^2} \right)^2 f_Q Y_u Y_u^\dagger Y_d f_d^2 Y_d Y_u Y_u^\dagger f_Q, \end{aligned} \quad (4.17)$$

and from the right:

$$\begin{aligned}
A_d &\equiv f_d \left( Y_d^{(\text{tot})} \right)^\dagger f_Q^2 Y_d^{(\text{tot})} f_d \\
&\sim f_d Y_d f_Q^2 Y_d f_d + \frac{4N_{\text{KK}}^2}{16\pi^2} f_d \left[ Y_d f_Q^2 Y_u Y_u^\dagger Y_d + Y_d Y_u Y_u^\dagger f_Q^2 Y_d \right] f_d \\
&+ \left( \frac{4N_{\text{KK}}^2}{16\pi^2} \right)^2 f_d Y_d Y_u Y_u^\dagger f_Q^2 Y_u Y_u^\dagger Y_d f_d,
\end{aligned} \tag{4.18}$$

The KK gluon coupling is diagonal but not universal in the original gauge eigenstate basis and is proportional to  $g_{s*} f_{q_i}^2$  for the  $q_i$  quark. The FCNC contributions are then proportional to the amount of rotation we need to perform in order to diagonalize (4.17-4.18). The terms linear in the loop factor have the same wave-function and RS-GIM structure as the tree-level terms in the generic RS model without alignment, so the only difference would be the suppression factor  $\epsilon_u$ . So the off-diagonal coupling of the LH down-type fields will be given by

$$g_{12}^L = g_{s*} \epsilon_u f_{Q_1} f_{Q_2}. \tag{4.19}$$

When considering a similar effect in the right handed fields, one also has to take into account another important fact: the flavor changing effects in the terms proportional to the square of the loop factor can also go through the third generation, and those have an additional enhancement factor of  $(f_{d_3}/f_{d_2})^2$  (for the LH fields) and  $(f_{Q_3}/f_{Q_2})^2$  (for the RH fields). Expressing these in terms of the physical quark masses and the Cabbibo angle, we find that for  $A_d$ , the term with the two-loop factor is actually larger than the term with the one loop-factor (while for  $A_Q$  they are roughly comparable). In fact since  $f_{Q_3}/f_{Q_2} \sim 1/\sin^2 \theta_c = 25$  this will be the leading term in the mass matrix, which implies that for the RH couplings one reverts back to the anarchic case - the misalignment in the RH sector is complete. So the RH couplings of the KK gluon will be unsuppressed compared to the standard RS GIM result

$$g_{12}^R = g_{s*} f_{d_1} f_{d_2}. \tag{4.20}$$

The effect of the KK gluon exchange is then  $g_{12}^L g_{12}^R / M_G^2$ , where as usual one converts to physical masses and mixing angles using the expressions of anarchic RS flavor (which are also applicable here), giving rise to

$$C_{4K} \sim \epsilon_u \frac{g_{s*}^2}{M_G^2} \frac{1}{Y_d^2} \frac{2m_d m_s}{v^2} \tag{4.21}$$

Thus naively one does not gain much (just one factor of  $\epsilon_u$ ) by going to alignment models. This is however not quite the case. We argue below that there are two ways using the alignment structure to further reduce the size of these corrections.

### Suppressing $\epsilon_K$ via enhanced $\langle y_d \rangle$

In order to numerically compare the bound on (4.21) to the most stringent bound on simple anarchic RS flavor from [9], we also need to take into account the fact that here the  $Y_d$

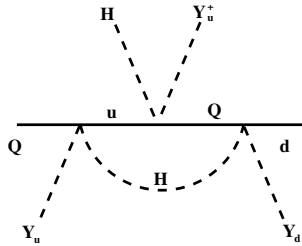


Figure 4: A 1-loop diagram on the IR brane responsible for the misalignment between the bulk and the brane.

coupling can be taken to be somewhat larger than the NDA bound in anarchic RS, since it is originating from a VEV of a bulk scalar. In this case, as we saw, raising the effective  $Y_d$  above the NDA value does not imply leaving the regime of validity of the effective theory, but rather some amount of tuning in the scalar potential setting the IR brane VEV. In the calculation of the bound on the KK gluon mass in [9], the  $Y_d$  coupling was fixed to  $Y_d \sim 3$ , and we have seen that  $C_{4K}$  in Eq.(4.21) scales as  $1/Y_d^2$ . The bound on the KK scale from this KK gluon mediated LLRR operator is about 10 TeV, if the running of the UV brane induced localized gluon kinetic term is taken into account (while it is about 20 TeV without that effect). If we want to reduce that bound to about 3 TeV (the lowest allowed by generic electroweak precision constraints), then one would need to raise  $Y_d$  to about

$$Y_d \gtrsim 3 \left( \frac{10 \text{ TeV}}{3 \text{ TeV}} \right) \sqrt{\epsilon_u} \sim 7.6 \left( \frac{N_{\text{KK}}}{3} \right). \quad (4.22)$$

Thus, raising the effective  $Y_d$  to about  $Y_d \geq 7.5$  (for  $N_{\text{KK}} = 3$ ) is sufficient (for  $N_{\text{KK}} = 2$ , one needs  $Y_d \geq 5.1$ ). This would imply that the  $\langle y_d \rangle$  VEV is about 2.5 times its NDA value, corresponding to a tuning of about 15% in the IR brane localized quartic. For  $N_{\text{KK}} = 2$ ,  $\langle y_d \rangle$  is just 1.1 times the NDA value, and the fine tuning is almost completely absent (just 75%). Note, that numerical results for aligned solutions with sufficiently large  $Y_d$  can be found both for the quadratic alignment (section 3.1) and for the linear alignment model (section 3.2). Also note that contributions to dipole operators from higher dimensional operators (such as EDMs,  $b \rightarrow s\gamma$  and  $\epsilon'/\epsilon_K$ , see [16] for a recent discussion) are suppressed by  $\epsilon_u$  and even further suppressed in the two-higgs doublet model described in the following section, and hence are probably consistent with the current bounds.

## Two-Higgs Doublet Model

Above we saw that the generic one-loop correction to the down Yukawa is quite large, so the final suppression is only proportional to  $\epsilon_u$  (and possibly further suppressed via enhanced down Yukawa couplings). One may try to suppress the misaligning terms further by introducing more structure. However, it is clear from the outset that there is no symmetry that can forbid the operator  $\bar{Q}Y_u^\dagger Y_u Y_d d H$  if the  $\bar{Q}Y_d d H$  is allowed, so this program can only postpone the issue to higher loops, which numerically might still be sufficient. One possible

way to realize this goal is if there were separate up and down type IR-localized Higgses: the  $H_u$  coupling to  $\bar{Q}Y_u u$  and the  $H_d$  to  $\bar{Q}Y_d d$ . Then the dangerous one-loop diagram would be absent. However, in the limit when there is no coupling between  $H_u$  and  $H_d$  there is also a new  $U(1)_{PQ}$  global symmetry which would give rise to a weak-scale axion. As usual this symmetry has to be broken, but the breaking scale (for example via a  $\mu^2 H_u H_d$  term) can be much lower than the KK scale  $1/R'$ , which will imply that the one-loop diagrams will be suppressed by  $(\mu/m_{KK})^2$ . If  $\mu \sim 100$  GeV, then we gain at least around  $1/100$  suppression and the dangerous one-loop diagram is rendered harmless. However, since there is no actual symmetry, the dangerous operator will indeed be generated at higher loops. For example, at 2 loops the diagram in Fig. 5 will be quartically divergent, and one can close up legs to get higher loop diagrams with higher power divergences generated. While the 2-loop diagram will have additional Higgs VEVs suppressing the contribution, one can close up those Higgs legs and obtain terms which have again  $\epsilon_u = \frac{Y_u^2 N_{KK}^2}{4\pi^2}$  as the expansion coefficient. While we have not attempted to study systematically the higher loop induced misaligning corrections, we can gain some idea of their magnitudes by estimating the magnitudes of the misaligning effects obtained by closing up legs on the diagram in Fig. 5. We find the following suppression factors in  $C_{4K}$  (for  $Y_d = 3$ ,  $Y_u = 1.6$ ,  $m_{KK} = 3$  TeV,  $\mu = 100$  GeV, taking into account that the Higgs VEVs in a 2HDM are smaller, and choosing the best case  $\tan \beta = 1$  for the numerical results):

#of loops	parametric suppression	$N_{KK} = 3$	$N_{KK} = 2$
1 – loop	$\epsilon_u \left( \frac{\mu}{m_{KK}} \right)^2$	$6 \cdot 10^{-4}$	$3 \cdot 10^{-4}$
2 – loop	$\frac{1}{\cos^2 \beta} (2\epsilon_u \frac{v}{m_{KK}} Y_d)^2$	0.16	0.03
3 – loop	$\frac{1}{\cos^2 \beta} \epsilon_u^3 \frac{Y_d^2}{Y_u^2}$	1.4	0.12
4 – loop	$\frac{1}{\cos^2 \beta} \epsilon_u^4 \frac{Y_d^4}{Y_u^4}$	0.31	0.01
5 – loop	$\frac{1}{\cos^2 \beta} \epsilon_u^5 \frac{Y_d^6}{Y_u^6}$	0.07	$10^{-3}$

We can see that the suppression is sufficient in all terms investigated if  $N_{KK} = 2$ , however for  $N_{KK} = 3$  the 3-loop contribution to the misaligning effects leading to  $C_{4K}$  is actually unsuppressed compared to the RS-GIM case.

### 4.3 Loop Induced Kinetic Mixings

The second potentially dangerous brane localized loops contribute to kinetic mixing terms for the left-handed ( $Q$ ) fields (at one loop) and right-handed fields (at two loops). The relevant diagrams are shown in Fig. 6.

The NDA estimate for the one-loop induced left-handed kinetic mixing term is

$$(\bar{Q} \not{D} Q)_{IR}^{(\text{loop})} \sim \bar{Q} \not{D} \frac{2N_{KK}}{16\pi^2} Y_u Y_u^\dagger Q. \quad (4.23)$$

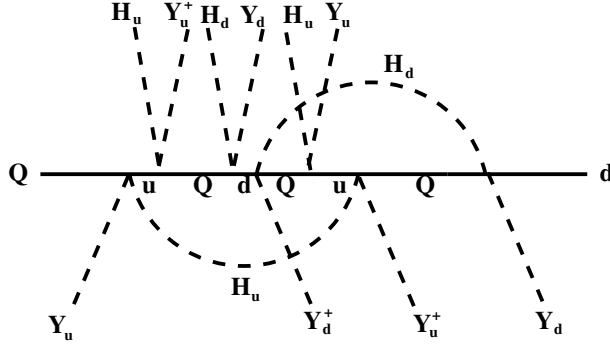


Figure 5: A 2-loop quartically divergent diagram contributing to the misalignment even in the 2HDM. By closing up additional Higgs and Yukawa scalar lines one can obtain higher loop (but more divergent) contributions.

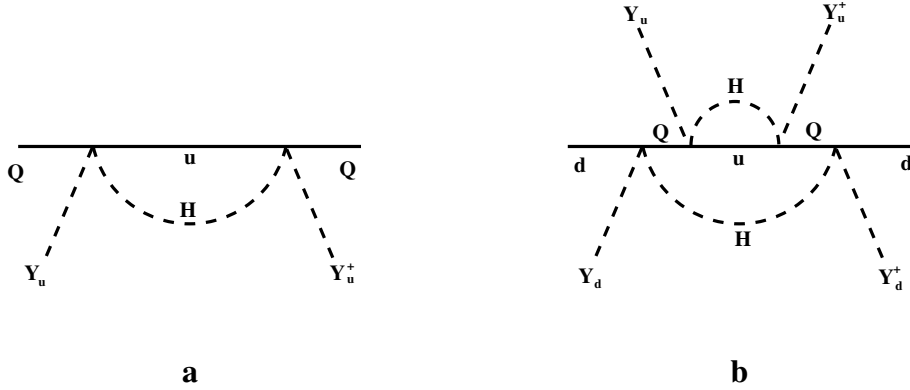


Figure 6: Corrections to the brane kinetic terms that will result in misaligned couplings, at one loop for  $Q$  in a and at two loops for  $d$  in b.

For the RH down quarks the NDA estimate for the two-loop contribution is

$$(\bar{d}\not{D}d)_{IR}^{(2\text{-loop})} \sim \bar{d}\not{D} \frac{4N_{\text{KK}}^3}{(16\pi^2)^2} Y_d Y_u Y_u^\dagger Y_d d. \quad (4.24)$$

The off-diagonal kinetic terms will directly contribute to the off-diagonal KK gluon coupling, since the kinetic term contains a covariant derivative including the KK gluon on the brane. Therefore the coefficient of the LR operator contributing to  $\epsilon_K$  (compared to generic magnitude for the anarchic case) is suppressed by the factor

$$g_{LLRR}/g_{RS1} \sim \frac{8N_{\text{KK}}^4}{(16\pi^2)^3} Y_d^2 Y_u^4 = \frac{\epsilon_u^2}{2} \frac{Y_d^2}{16\pi^2} \quad (4.25)$$

which for  $N_{\text{KK}} = 3$  is already a suppression by  $1/35$ , thus putting it safely within the bounds.

## 4.4 UV Brane Kinetic Terms

Next we consider the effects of flavor violation on the UV brane. The analysis here is quite general, and does not rely on the specific alignment structure of the bulk mass parameters. We will allow arbitrary breaking of the flavor symmetries, and assume that the main effect of this is to induce a kinetic mixing matrix for the quarks on the UV brane. The only assumption we are making is that the localized kinetic terms do not dominate over the bulk kinetic terms, but could be an  $\mathcal{O}(1)$  correction (otherwise they would significantly change the mass hierarchy that is assumed to come from the wave function overlaps). This is equivalent to the assumption that the elementary/composite ratio of the SM fermions is determined by the bulk  $c$ 's, which a priori can be modified by large localized kinetic terms (which is the approach taken in [22]). Therefore, we write

$$\mathcal{L}_{UV} = K_{ij} R \bar{\psi}_L^j \gamma_\mu D^\mu \psi_L^i \quad (4.26)$$

where  $K$  is a dimensionless kinetic mixing matrix. The kinetic terms of the zero modes can be diagonalized via a hermitian rotation  $H = U N^{-\frac{1}{2}} U^\dagger$ , such that

$$1 = H(1 + \tilde{f} K \tilde{f}) H, \quad (4.27)$$

where the  $\tilde{f}$ 's are the dimensionless fermion wave functions on the UV brane

$$\psi^i(z = R) = \frac{1}{\sqrt{R}} \tilde{f}^i. \quad (4.28)$$

We can express  $H$  as

$$H = \left(1 + \tilde{f} K \tilde{f}\right)^{-\frac{1}{2}}. \quad (4.29)$$

Now we need to rotate the KK gluon couplings using  $H$ . Using the approximate expression for the overlap integral of the KK gluon with the bulk fermions, and adding the contributions

from the UV induced kinetic term (which also contains a coupling to the gluon) and using the fact that the KK gluon wave-function on the UV brane is  $1/\log(R'/R)$ , we obtain

$$\begin{aligned} g_{ffG} &\approx g_{s*} H \left[ -\frac{1}{\log(R'/R)} + f_c^2 \gamma(c) - \frac{1}{\log(R'/R)} \tilde{f} K \tilde{f} \right] H \\ &= g_{s*} \left[ -\frac{1}{\log(R'/R)} + H f_c^2 \gamma(c) H \right], \end{aligned} \quad (4.30)$$

We can see that there is an automatic RS-GIM suppression for the UV effects of the UV brane kinetic terms too! This is not so surprising: both the KK gluon and the SM fermions can be thought of as mixtures of elementary and composite fields. The elementary KK gluon is the KK gluon at the UV brane, which must have flavor universal couplings. So all the effects of the kinetic terms must go through the composite part of the KK gluon (and the SM fermions), which will necessarily imply that the RS-GIM protection, which is what we have explicitly found here. We can go a step further and estimate bounds on the off-diagonal terms of  $K$ , assuming it is a subdominant effect compared to 1. In that case

$$H \approx 1 - \frac{1}{2} \tilde{f} K \tilde{f}. \quad (4.31)$$

The ratio of the LR operator  $C_4^K$  compared to the standard anarchic case will be

$$g_{LLRR}/g_{RS1} \sim \frac{1}{4} K_{12}^L K_{12}^R \frac{m_s}{m_d}. \quad (4.32)$$

Again requiring that a 3 TeV KK gluon is satisfying the bound gives the constraint

$$K_{ij} \lesssim 0.1. \quad (4.33)$$

The relative contribution of third generation effects at each of the vertices is given by

$$\frac{K_{13} K_{32}}{2K_{12}} \tilde{f}_{c_3}^2 \frac{f_{c_3}^2}{f_{c_2}^2} = \frac{(1 - 2c_3)(R/R')^{1-2c_3}}{1 - (R/R')^{1-2c_3}} \begin{cases} \frac{K_{13}^L K_{32}^L}{2K_{12}^L} \frac{1}{\lambda^4} \\ \frac{K_{13}^R K_{32}^R}{2K_{12}^R} \left( \frac{m_b \lambda^2}{m_s} \right)^2 \end{cases} \quad (4.34)$$

we see that third generation effects are exponentially suppressed for  $c_3 < 1/2$ . Even for  $c_3 > 1/2$ , the extra power of  $K$  renders third generation contributions subdominant (we have assumed that the entries of  $K$  satisfy the bound in (4.33)). We learn that if  $K$  were loop generated, then the extra loop factor would make these bounds satisfied easily. Of course generically there is no reason for the  $K$  to be loop suppressed. In the alignment model, one can of course eliminate all the FCNC's from the UV brane by postulating that the only source breaking the flavor symmetries is again proportional to the  $Y_d$  field itself. In this case, the boundary kinetic term is itself aligned with the bulk  $c_{Q,d}$  and the Yukawa coupling, and no flavor violation would occur.

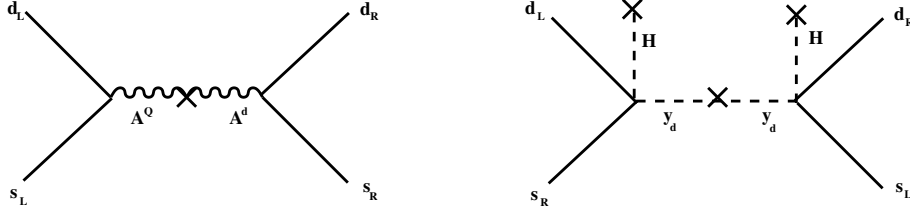


Figure 7: Diagrams contributing to LR down-type FCNCs. These interactions are induced by exchanging either flavor gauge bosons (left) or flavor scalars (right). The insertions are flavor violating mass terms, which in the alignment models are proportional to  $Y_d$ .

## 4.5 More Contributions to FCNC

In our setup we have at least two additional sets of fields that could potentially lead to flavor violating effects:

- the bulk scalar  $y_d$  which is responsible for the alignment
- bulk flavor gauge bosons (which are the gauge fields for the bulk flavor symmetries  $SU(3)_Q \times SU(3)_d$ ).

The couplings of these fields (before  $y_d$  obtains a VEV) are flavor invariant, since the flavor symmetries are assumed to be gauge symmetries. Thus the only source of flavor violation can be due to the VEV of the scalar field  $y_d$  which is generating flavor violating masses for the flavor gauge bosons, the down Yukawa coupling, and potentially also flavor violating masses for the dynamical scalar KK modes. Exchange of the heavy (TeV-scale) gauge and scalar modes can in principle give rise to the most dangerous flavor changing LR four-Fermi operators in the down sector, *e.g.* from diagrams of the form presented in Fig. 7. However, for the alignment models presented in this paper the only source of flavor violation in the bulk is due to the VEV  $\langle y_d \rangle \sim Y_d$ , so the alignment extends to all the effects in these sectors too: the flavor gauge bosons and the KK modes of  $y_d$  will be flavor diagonal in the basis where the down-quark masses are diagonalized, thus no additional FCNC's will be introduced. In the general case with multiple bulk scalars  $\propto Y_u, Y_d$  there will be (RS-GIM suppressed) effects comparable to the size of the KK-gluon exchange generated (depending of course on the size of the 5D gauge flavor gauge coupling). We leave for future work the detailed study of the properties of these fields and their flavor violating effects in the general case without alignment.

## 5 Conclusions

In this work, we have shown that warped extra dimensions may provide the solution both to the hierarchy problem and the flavor puzzle, while being accessible at the LHC. This can be achieved by invoking flavor alignment between the down-type 5D Yukawas and the bulk fermion masses. Such alignment reduces the flavor problems which are usually present in low-scale models of new physics. In order to have 5D flavor alignment in the bulk, we assume



a flavor breaking source at the UV brane, and promote the Yukawa couplings to bulk scalar fields, so that they can shine the flavor violation down the bulk. We have discussed two models with alignment which are improved versions of the generic 5DMFV model. Potential problems of this idea - namely loop induced misalignment, flavor violation from UV brane kinetic terms, and flavor violation induced by the new bosons - were shown not to spoil the alignment. Our framework is fairly predictive in its flavor structure which should allow to test it, possibly even in the near future. One of the most robust prediction is that the up sector is anarchical. This should induce CPV in the  $D$  system which is just around the corner. Furthermore, top FCNC contribution via  $t_R \rightarrow cZ$  is around the LHC reach [37]. However, in models where  $t_R$  transforms trivially under the custodial symmetry [32] to protect  $Z \rightarrow b\bar{b}$ , this effect is absent. It is interesting that the same selection of trivial RH up type quark representation also eliminates (to leading order) Higgs mediated FCNCs [17] which correlates the smallness of the RS contributions to  $\epsilon_K$  with the absence of top flavor violation. In addition we note that the alignment is more effective in suppressing flavor violation between the third generation and the light ones (the enhancement which is proportional to the ratio of the third to second fermion compositeness fraction is absent). Thus, our model actually predicts that these transitions are suppressed by  $\epsilon_u \sim 0.5$  compared to the anarchic case. It suggests a discovery of CP violation in  $D^0 - \bar{D}^0$  mixing and a small deviation from the SM predictions in the the down sector might provide a way to verify or exclude our scenario.

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## A Fermion zero mode with a bulk mass from shining

We show how the wave function of the fermion zero mode is affected by contributions from a bulk scalar field. The action for a bulk scalar field  $\phi$  with a Yukawa v.e.v. on the UV brane is

$$\begin{aligned}
S = & \int d^5x \left( \frac{R}{z} \right)^5 \left\{ \left( \frac{z}{R} \right)^2 \frac{1}{2} [\partial_\mu \phi \partial^\mu \phi - (\partial_z \phi)^2] - \frac{1}{2} m^2 \phi^2 \right\} \\
& - \int d^4x \left( \frac{R}{z} \right)^4 \frac{\lambda}{2R} (\phi - Y)^2 \Big|_R
\end{aligned} \tag{A.1}$$

With the boundary conditions

$$0 = \partial_z \phi \Big|_{R'}, \quad 0 = \left( \partial_z \phi - \frac{\lambda}{R} (\phi - Y) \right) \Big|_R, \tag{A.2}$$

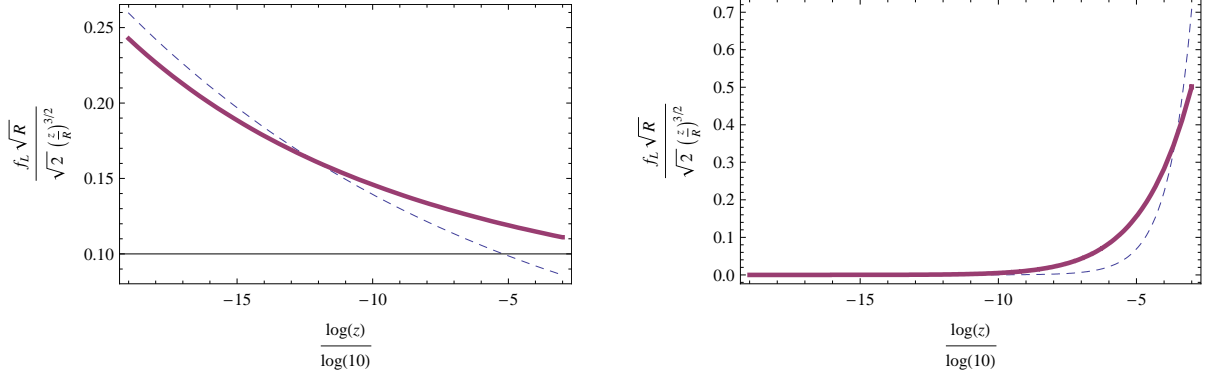


Figure 8: Normalized zero modes. Dashed: standard bulk mass term with  $c = 1/2 + \beta$  and no scalar contribution, solid line:  $c = 1/2$  and additional scalar vev  $\phi \sim \beta \left(\frac{R}{z}\right)^{2\epsilon}$ . Left Panel:  $\beta = 0.05$ , Right panel:  $\beta = -0.506$  ( $\epsilon = 0.01$  corresponding to  $MR = 1/2$ ). One can read of the value  $f(c)$  of the wave function on the IR brane.

we get

$$\phi(z) = Y \left( \frac{R}{z} \right)^\epsilon. \quad (\text{A.3})$$

where we have set  $\epsilon = \sqrt{(mR)^2 + 4} - 2$  and assumed  $\epsilon \ll 1 \ll \lambda$ . The bulk mass term of the fermion receives additional contributions from shining

$$\int d^5x \left( \frac{R}{z} \right)^5 M \psi \chi \rightarrow \int d^5x \left( \frac{R}{z} \right)^5 \psi \chi \left( M + \alpha \frac{\phi^\dagger \phi}{\Lambda^2} + \dots \right) \quad (\text{A.4})$$

$$\rightarrow \int d^5x \left( \frac{R}{z} \right)^5 \psi \chi \left( M + \alpha \frac{Y^\dagger Y}{\Lambda^2} \left( \frac{R}{z} \right)^{2\epsilon} + \dots \right) \quad (\text{A.5})$$

Setting  $\beta = \alpha \frac{Y^\dagger Y}{R \Lambda^2}$ , the EOM for the left-handed zero mode becomes

$$\left( \partial_z - \frac{2-c}{z} \right) f_L(z) + \frac{1}{z} \beta \left( \frac{z}{R} \right)^{-2\epsilon} f_L(z) = 0, \quad (\text{A.6})$$

where we have set  $c = MR$ . The solution is given by

$$f_L(z) = N(\epsilon) z^{2-c} \exp \left[ \frac{\beta}{2\epsilon} \left( \frac{R}{z} \right)^{2\epsilon} \right] = \tilde{N}(\epsilon) z^{2-c-\beta} \left( 1 + \beta \epsilon \left( \ln \frac{R}{z} \right)^2 + \dots \right) \quad (\text{A.7})$$

In the limit  $\epsilon \rightarrow 0$ , we find the usual form of the zero mode with a bulk mass  $c + \beta$ . The leading correction to the shape of the wave-function is given by the second term and we show two examples in Fig. 8.

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